

# Instant nonthermal leptogenesis

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We propose an economical model of nonthermal leptogenesis following inflation during “instant” preheating. The model involves only the inflaton field, the standard model Higgs, and the heavy “right-handed” neutrino.

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## I. INTRODUCTION

Leptogenesis [1] is an attractive scenario to account for the observed matter–antimatter asymmetry of the universe. In the scenario, a lepton asymmetry is generated by the decay of massive right-handed (Majorana) neutrinos,  $N$ , which are responsible for the (small) masses of left-handed neutrinos via the see-saw mechanism [2]. The lepton asymmetry is then translated to a baryon asymmetry by sphaleron processes [3] around the electroweak era. Leptogenesis is analogous to GUT baryogenesis [4] in that a massive particle decays with CP violation, and Boltzmann equations may be employed to track the time evolution of particles and the asymmetry [5]. The massive particles (in the case of leptogenesis, the  $N$ s) must be created after inflation, either nonthermally or thermally during reheating, or thermally during the radiation-dominated era. Thermal leptogenesis [6, 7, 8, 9] has been widely studied, but a careful analysis must be done in order to account for all the thermal and perturbative reactions (see Ref. [8] and Ref. [9]). Nonthermal leptogenesis is an attractive alternative [10, 11]. Here, we propose a new model of nonthermal leptogenesis involving instant preheating.

Our model assumes hybrid inflation [12], in which inflation is terminated by an abrupt transition in the properties of the scalar-field potential dominating the energy density during inflation. Such an abrupt change is often modeled as being triggered the action of a second “waterfall” field, causing the effective scalar-field to roll into another dimension in the scalar-field landscape. An important feature of hybrid inflation is that the properties of the scalar-field potential (*e.g.*, the mass) during reheating may be quite different than the properties of the scalar-field potential during inflation.<sup>1</sup>

Our model assumes that the scalar-field energy is extracted and thermalized by instant preheating. In preheating [13], particles are produced when the inflaton passes through a nonadiabatic phase around the minimum of the inflaton potential. In “instant” preheating [14], the inflaton is strongly coupled to a particle whose mass depends on the value of the inflaton field. This particle can be either a boson [13] or a fermion [11]; in our model we will assume it is a boson. As the inflaton oscillates, the coupling of the inflaton to the produced particle results in an increasing mass of the produced particle. As the mass of the produced particle increases, its decay rate will also increase, and decay channels disallowed when the produced particle is at the minimum of its potential may open. We will take advantage of both these properties.

In our model we will assume that the inflaton couples to the standard-model Higgs boson,  $h$ , associated with electroweak symmetry breaking. We will also assume that, as expected,  $h$  couples to the  $N$ . Normally the mass of the Higgs,  $m_h$ , is much, much less than the mass of the  $N$ ,  $m_N$ . However, during instant preheating this need not be

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<sup>1</sup> Although the scalar-field landscape may be quite complicated, involving several degrees of freedom, we will refer to the scalar field during preheating as the inflaton, although as remarked, the mass of the “inflaton” after inflation may be different than the mass of the inflaton during inflation.

the case, and the Higgs may decay directly into  $N$ , producing a lepton asymmetry. Later when the inflaton is close to its minimum, the produced  $N$ s become heavier than the Higgs, and they will decay back to the Higgs.

Let us elaborate on this picture. We will assume that the temperature,  $T$ , is always less than the mass of the  $N$  throughout the entire preheating stage. The mass of the Higgs will be determined by its coupling to the inflaton  $\phi$ :  $m_h \propto |\phi|$ . For sufficiently large values of  $|\phi|$  during the inflaton oscillations,  $m_h$  will be larger than  $m_N$ . We will denote the absolute value of  $\phi$  when  $m_h(\phi) = m_N$  as  $\phi_c$ .

It is useful to imagine a single oscillation of the inflaton field, in particular the first oscillation. As  $\phi$  passes near its minimum,  $h$  is effectively massless, and a burst of  $h$ s are created. The  $h$ s will decay to any kinematically allowed final states. Because of the large  $h$ -top-quark coupling, the decay is predominately into top quarks.  $h \rightarrow N$  becomes kinematically allowed when  $|\phi|$  becomes larger than  $\phi_c$ . Therefore, efficient lepton number production happens when  $\phi_c$  is close to the minimum so  $h \rightarrow N$  process takes place before all the  $h$  decays thermally into top quarks. In the case of hierarchical  $N$ s, where  $m_{N_1} = g|\phi_c| \ll m_{N_2} = g|\phi_2| \ll m_{N_3} = g|\phi_3|$ ,  $h$ s decay while  $|\phi| \ll |\phi_{2,3}|$  due to the large  $h$ -top-quark coupling. This process is nonthermal, as  $m_h > T$  at this time. Eventually  $\phi$  reaches a maximum point  $\phi_0^{\max}$  and rolls back down. The decay of the  $h$  continues until  $\phi < \phi_c$ . At this stage,  $N \rightarrow h$  decay happens, and  $h$ s continue to decay into fermions. A lepton asymmetry is generated by both  $h \rightarrow N$  and  $N \rightarrow h$  decays. Another burst of  $h$ s are produced as  $\phi$  passes again through the nonadiabatic phase at the origin, and the same events occur on the other side of the potential. Since  $h$  decays very rapidly, a negligible amount of  $h$ s remain when  $\phi$  repasses through the nonadiabatic regime to produce more  $h$ s. This eliminates the influence of the old  $h$ s with  $\phi$  during production of new  $h$ s, and the backreaction of Higgs in the nonadiabatic region need not be considered. The production and decay of  $h$ s siphon away energy from  $\phi$ , and  $\phi^{\max}$  decreases for each oscillation. A schematic diagram of the regions of the potential in instant preheating is shown in Fig. 1.

There have been many models of leptogenesis. A hallmark of our model is the economy of fields. The only undiscovered fields are the inflaton,  $\phi$ , the standard model Higgs,  $h$ , and the right-handed neutrino,  $N$ . There are very good reasons for suspecting that all exist! The only unfamiliar aspect of our model is the strong coupling of the inflaton field to the Higgs field. While there is no reason to preclude such a coupling, it would be very interesting to find particle-physics models with a motivation for the coupling.

In the next section we present the Lagrangian used in our calculation, we parameterize the nonadiabatic creation of particles in preheating, and we discuss the decay rates and CP violation parameters. In Section III we present and solve the Boltzmann equations used in the calculation, presenting the main results in the form of figures. The final section contains our conclusions.

## II. INSTANT PREHEATING AND THE SEE-SAW MECHANISM

Inflation ends in the hybrid model when  $\phi$  meets a “waterfall” potential in another direction of the scalar field landscape. The  $\phi$  promptly falls into this potential which is responsible for preheating. Hence,  $\phi$  does not carry restrictions on potential parameters (such as the mass) deduced from present cosmological observations. For instance, the mass of the  $\phi$  during the preheating process may be more massive than the mass of the  $\phi$  during inflation.

We assume  $\phi$  is coupled to the standard model Higgs  $h$ , with interaction Lagrangian [14] of the preheat field given by

$$\mathcal{L}_{\text{preheat}} = -\frac{1}{2}g^2\phi^2h^2, \quad (1)$$

where  $g$  is the coupling constant. Ignoring its electroweak-scale mass,  $m_h = g|\phi|$ . We define  $\phi_c \equiv g/m_N$ . Thus, depending on the initial condition of the  $\phi$  field,  $m_h$  may become larger or smaller than  $m_N$  as  $\phi$  oscillates about the minimum of its potential.

The inflaton-Higgs coupling leads to a potential of  $\phi$  about the minimum in the form

$$V(\phi) = \frac{1}{2}\mu^2\phi^2 + \frac{1}{2}g^2\phi^2h^2, \quad (2)$$

where  $\mu$  is the  $\phi$  mass. The Hartree approximation will be used to take the average value of  $h^2$  [13], where we approximate  $h^2 \approx \langle h^2 \rangle$ . Of course  $h^2$  is formally infinite, but becomes  $2n_h/m_h$  after renormalization. The equation of motion of  $\phi$  becomes

$$\ddot{\phi} + 3H\dot{\phi} + \mu^2\phi + 2gn_h\phi/|\phi| = 0, \quad (3)$$

where the dot stands for the time derivative,  $H \equiv \dot{a}/a$  is the Hubble expansion parameter, and  $a$  is the scale factor.

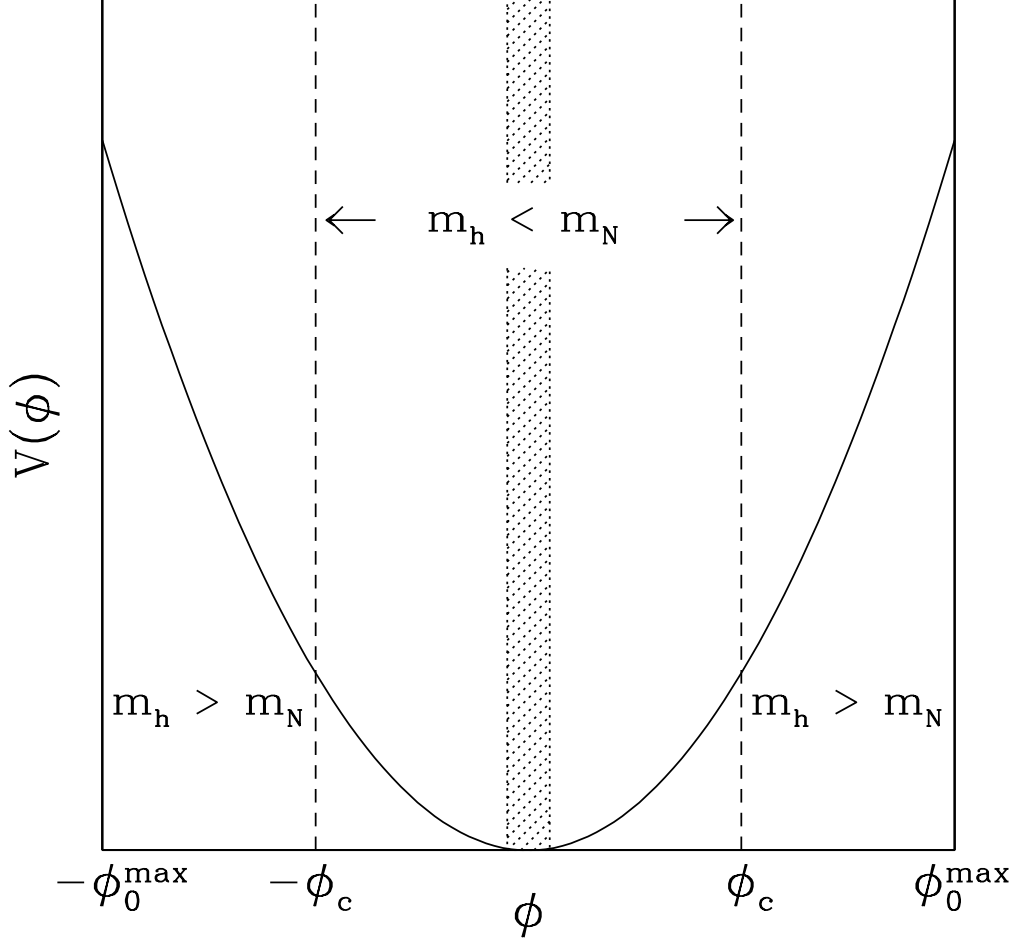


FIG. 1: A schematic diagram of regions in the inflaton potential during instant preheating. The shaded column around the minimum illustrates the nonadiabatic region where  $h$ s are created. In regions of  $|\phi| > \phi_c$ ,  $h \rightarrow N$  decay occurs. In regions of  $|\phi| < \phi_c$  (modulo the nonadiabatic region),  $N \rightarrow h$  decay occurs. The regions are not drawn to scale.

The  $h$ s are created when  $\phi$  goes through a nonadiabatic phase, which occurs near the minimum of the potential [13]. This phase is very short and can be treated as instantaneous. The number density of  $h$ s created in the nonadiabatic phase is [13, 14]

$$n_h(0) = \frac{1}{2\pi^2} \int_0^\infty dk k^2 n_k = \frac{1}{2\pi^2} \int_0^\infty dk k^2 e^{-\pi k^2 / g |\dot{\phi}_0|} = \frac{(g \dot{\phi}_0)^{3/2}}{8\pi^3}, \quad (4)$$

where  $\dot{\phi}_0$  is the initial time derivative of  $\phi$ , and  $k$  is the momentum. A large coupling constant  $g \sim 1$  enables a quick and effective thermalization of the universe within a few oscillations of  $\phi$ .

The see-saw mechanism Lagrangian for three families with Majorana neutrino masses  $m_{N_i}$  ( $i = 1, 2, 3$ ) and Yukawa couplings  $Y_{ij}^\nu$  to the Higgs and light neutrinos  $l$  is given by

$$\mathcal{L}_{\text{see-saw}} = \frac{m_{N_i}}{2} N_i^2 + Y_{ij}^\nu l_i N_j h. \quad (5)$$

The left-handed light neutrino masses are  $m_\nu = -(vY^\nu)^T m_N^{-1} (vY^\nu)$ , where  $v = 247$  GeV is the Higgs vacuum expectation value.  $\mathcal{L}_{\text{see-saw}}$  also generates a dimension-5 effective operator which causes CP violation among the leptons. Throughout the paper we consider the case of hierarchical Majorana neutrinos,  $m_{N_1} \ll m_{N_2} \ll m_{N_3}$ . Hence the CP asymmetry is created during  $|\phi_c| \leq |\phi| \ll |\phi_{2,3}|$ , as the  $h \rightarrow Nl$  interaction competes against the dominant background interaction of  $h \rightarrow f\bar{f}$ . This mass hierarchy allows us to consider only interactions involving  $N_1$ ; hence we will drop the family subscript unless distinction between  $N_1$  and  $N_{2,3}$  is required.

The decay processes  $h \rightarrow Nl(\bar{l})$  or  $N \rightarrow hl(\bar{l})$  give rise to a lepton asymmetry. Both processes are possible in our model: the former process happening when  $m_h > m_N$  (*i.e.*, when  $|\phi| > \phi_c$ ), and the latter happening for  $m_h < m_N$  (*i.e.*, for  $|\phi| < \phi_c$ ). The CP parameters in these interactions,  $\epsilon_h$  and  $\epsilon_N$ , respectively, are defined as

$$\epsilon_h \equiv \frac{\Gamma_{h \rightarrow Nl} - \Gamma_{h \rightarrow N\bar{l}}}{\Gamma_{h \rightarrow Nl} + \Gamma_{h \rightarrow N\bar{l}}}; \quad \epsilon_N \equiv \frac{\Gamma_{N \rightarrow hl} - \Gamma_{N \rightarrow h\bar{l}}}{\Gamma_{N \rightarrow hl} + \Gamma_{N \rightarrow h\bar{l}}}, \quad (6)$$

where the subscripts of the decay width  $\Gamma$  denotes the decay process concerned. The possible combination of these CP violating processes are

$$h \longrightarrow \begin{cases} Nl \rightarrow \begin{cases} hll \\ h\bar{l}\bar{l} \end{cases} \\ N\bar{l} \rightarrow \begin{cases} h\bar{l}\bar{l} \\ hll \end{cases} \end{cases}. \quad (7)$$

The second and third final states have zero lepton number, but the first and fourth final states give rise to lepton asymmetry. The total CP asymmetry  $\epsilon_{\text{tot}}$  is expressed as

$$\epsilon_{\text{tot}} \equiv \left( \frac{\Gamma_{h \rightarrow Nl}}{\Gamma_{h \rightarrow Nl} + \Gamma_{h \rightarrow N\bar{l}}} \frac{\Gamma_{N \rightarrow hl}}{\Gamma_{N \rightarrow hl} + \Gamma_{N \rightarrow h\bar{l}}} - \frac{\Gamma_{h \rightarrow N\bar{l}}}{\Gamma_{h \rightarrow Nl} + \Gamma_{h \rightarrow N\bar{l}}} \frac{\Gamma_{N \rightarrow h\bar{l}}}{\Gamma_{N \rightarrow hl} + \Gamma_{N \rightarrow h\bar{l}}} \right) = \frac{1}{2} (\epsilon_h + \epsilon_N). \quad (8)$$

The explicit expression of this CP parameter  $\epsilon$  at one-loop is [8]

$$\epsilon = -2 \sum_{i \neq 1} \frac{\text{Im}[(Y^\dagger Y)_{1i}^2] \text{Im}[I_0^* I_1]}{(Y^\dagger Y)_{11} 2P \cdot P_l}, \quad (9)$$

for both  $h$  and  $N$  decay.  $I_0$  and  $I_1$  are the tree- and one-loop diagrams, respectively. In Eq. (9),  $P$  is the four-momentum for either  $h$  or  $N$ , and  $P_l$  is the four-momentum of  $l$ . One of the condition of our leptogenesis model is the nonthermal production of  $N$ . All processes involving CP violation are nonthermal and a zero-temperature expression of  $\epsilon$  is required. Because of the mass hierarchy,  $N_{2,3}$  are not on-shell in the one-loop diagram. The  $h$  and  $l$  are considered massless ( $m_{h,l} \ll m_N$ ) in calculating  $\epsilon_N$ , and  $N$  and  $l$  are considered massless ( $m_{N,l} \ll m_h$ ) for the calculation of  $\epsilon_h$ . Therefore the expression of the above equation is equivalent for both  $\epsilon_h$  and  $\epsilon_N$ , and thus  $\epsilon_h = \epsilon_N$ . We look into the case of  $\epsilon_N$  for the  $N \rightarrow h$  decay in the following as an example.

In a hierarchical  $N$ -family structure ( $m_{N_1} \ll m_{N_{2,3}}$ ), the zero-temperature expression of Eq. (9) becomes

$$\epsilon_N = \frac{1}{8\pi} \sum_{i \neq 1} \frac{\text{Im}[(Y^\dagger Y)_{1i}^2]}{(Y^\dagger Y)_{11}} f\left(\frac{m_{N_i}^2}{m_{N_1}^2}\right), \quad f(x) = \sqrt{x} \left[ \frac{x-2}{x-1} - (1+x) \ln\left(\frac{1+x}{x}\right) \right]. \quad (10)$$

A convenient parameter to use is the effective neutrino mass [7, 8]

$$\tilde{m}_1 \equiv (Y_\nu^\dagger Y_\nu)_{11} \frac{v^2}{m_N}, \quad (11)$$

which is the contribution to the neutrino mass mediated by  $N_1$ . A way to understand  $\tilde{m}_1$  is to use an example of a hierarchical lefthanded neutrino spectrum of  $m_1 \ll m_2 = m_{\text{sun}} \ll m_3 = m_{\text{atm}}$  where  $m_{\text{sun}}$  and  $m_{\text{atm}}$  are the deduced solar and atmospheric neutrino mass from neutrino oscillations. If  $N_1$  gives rise to the atmospheric mass splitting then  $\tilde{m}_1 = m_{\text{atm}}$ ; if  $N_1$  causes the solar mass splitting then  $\tilde{m}_1 \gtrsim m_{\text{sun}}$ .

The explicit expression of  $\epsilon_N$  is [15]

$$|\epsilon_N| \leq \frac{3}{16\pi} \frac{m_N (m_3 - m_1)}{v^2} \times \begin{cases} 1 - m_1/\tilde{m}_1 & \text{if } m_1 \ll m_3 \\ \sqrt{1 - m_1^2/\tilde{m}_1^2} & \text{if } m_1 \simeq m_3 \end{cases}. \quad (12)$$

$m_3$ , has not been measured. We assume  $m_3 = \max(\tilde{m}_1, m_{\text{atm}}) = 0.05$  eV. Eq. (12) is maximal when  $m_1 = 0$ ;

$$|\epsilon_N^{\text{max}}| = \frac{3}{16\pi} \frac{m_{N_1} m_3}{v^2}. \quad (13)$$

The above expression holds for  $|\epsilon_h^{\text{max}}|$  as well.

The  $h$  decay thermalizes the universe. Before electroweak symmetry breaking in the standard model, only the fermion decay channel is allowed. The decay width of  $h \rightarrow f\bar{f}$  is given by

$$\Gamma_{h \rightarrow f\bar{f}} = 2 \frac{m_f^2}{v^2} m_h, \quad (14)$$

where  $m_f$  is the mass of the fermion. The decay widths of  $h \rightarrow Nl$  and  $h \rightarrow N\bar{l}$  differs only by the CP asymmetry. When calculating the Boltzmann equations, the CP asymmetry is factored out as  $\epsilon$  and an identical expression for  $\Gamma_{h \rightarrow Nl}$  and  $\Gamma_{h \rightarrow N\bar{l}}$  are used. In these cases,  $l\bar{l}$  will be dropped from the reaction expressions. The same holds for  $N \rightarrow hl$  and  $N \rightarrow h\bar{l}$ . The decay widths for these are

$$\Gamma_{h \rightarrow N} = (Y^\dagger Y)_{11} \frac{m_h}{8\pi}, \quad (15)$$

$$\Gamma_{N \rightarrow h} = (Y^\dagger Y)_{11} \frac{m_N}{8\pi}. \quad (16)$$

### III. LEPTOGENESIS

We study the time evolution of the  $hs$ ,  $Ns$ , and the lepton asymmetry by means of the Boltzmann equations. Along with the equation of motion for  $\phi$  in Eq. (3), the following set of Boltzmann equations are used:

$$\dot{n}_h + 3Hn_h + \Gamma_{h \rightarrow f\bar{f}}(n_h - n_h^{eq}) + \Gamma_{h \rightarrow N}(n_h - n_h^{eq}) - \Gamma_{N \rightarrow h}(n_N - n_N^{eq}) = 0, \quad (17)$$

$$\dot{n}_N + 3Hn_N + \Gamma_{N \rightarrow h}(n_N - n_N^{eq}) - \Gamma_{h \rightarrow N}(n_h - n_h^{eq}) = 0, \quad (18)$$

$$\dot{n}_L + 3Hn_L - \frac{\epsilon_h}{2}\Gamma_{h \rightarrow N}(n_h - n_h^{eq}) - \frac{\epsilon_N}{2}\Gamma_{N \rightarrow h}(n_N - n_N^{eq}) = 0, \quad (19)$$

$$\dot{\rho}_R + 4H\rho_R - \Gamma_{h \rightarrow f\bar{f}}(n_h - n_h^{eq}) - \Gamma_{h \rightarrow N}m_h(n_h - n_h^{eq}) - \Gamma_{N \rightarrow h}(n_N - n_N^{eq})m_N = 0. \quad (20)$$

Here,  $n_N$  is the number density of  $N$ ,  $n_L \equiv n_l - n_{\bar{l}}$  is the lepton number density, and  $\rho_R$  is the radiation energy density. It is understood that  $\Gamma_{h \rightarrow N}$  occurs when  $m_h > m_N$ , and  $\Gamma_{N \rightarrow h}$  occurs when  $m_h < m_N$ . The expansion rate  $H$  is

$$H^2 = \frac{8\pi}{3M_{Pl}^2} \left( \frac{1}{2}\dot{\phi}^2 + V(\phi) + \rho_h + \rho_N + \rho_R \right), \quad (21)$$

where  $M_{Pl}$  is the Planck mass, and  $\rho_h$  and  $\rho_N$  the  $h$  and  $N$  energy densities density.

There is an implicit assumption of a rapid thermalization from  $\phi$  and  $h$  decay products. The  $h$  decay products include fermions (Eq. (14)), where the dominant channel are decay into top quarks, which have large cross section and will result in rapid thermalization. Indeed, instant preheating relies on rapid thermalization from the large coupling of  $\phi^2$ - $h^2$  and fast decay of  $hs$  [14].

The  $\Delta L = 2$  off-shell scattering is not included in Eq. (19), but as its decay rate is very small compared to the on-shell scattering [8], we neglect this effect.

The  $hs$  are partially thermal throughout preheating. They are nonthermal when  $|\phi| > \phi_c$ , *i.e.*,  $m_h > m_N$ . They may be thermal when  $|\phi| < \phi_c$ , depending on the thermalization rate, and its effect must be taken into account. The equilibrium number density of  $h$ , using Maxwell-Boltzmann statistics, is

$$n_h^{eq} = \frac{T^3}{2\pi} \left( \frac{m_h}{T} \right)^2 K_2(m_h/T), \quad (22)$$

where  $K_2$  is the modified Bessel function of the second kind. The  $Ns$  are nonthermal throughout preheating. Consequently  $n_N^{eq} \gg n_N$ , and  $n_N^{eq}$  terms are neglected in the above Boltzmann equations.

We now express Eqs. (3) and (17)-(20) in terms of dimensionless variables and convert time derivatives to derivatives with respect to  $x \equiv \mu t$ . The dimensionless variables are

$$\begin{aligned} \varphi &\equiv \frac{\phi}{\phi_0^{\max}}; & R &\equiv \frac{\rho_R}{\phi_0^2/2} A^4; & \mathcal{H} &\equiv \frac{H}{\mu}; & A &\equiv \frac{a}{a_0}; \\ N_h^{eq} &\equiv \frac{n_h^{eq}}{n_h(0)} A^3; & N_h &\equiv \frac{n_h}{n_h(0)} A^3; & N_N &\equiv \frac{n_N}{n_h(0)} A^3; & N_L &\equiv \frac{n_L}{n_h(0)} A^3. \end{aligned} \quad (23)$$

Here,  $a_0$  is the initial scale factor. The dimensionless equation of motion of  $\phi$  and the Boltzmann equations become

$$\varphi'' + 3\mathcal{H}\varphi' + \varphi + \beta \left( \frac{N_h}{A^3} - \frac{N_h^{eq}}{A^3} \right) = 0, \quad (24)$$

$$N_h' - \frac{\Gamma_{h \rightarrow f\bar{f}}}{\mu} (N_h - N_h^{eq}) + \frac{\Gamma_{h \rightarrow N}}{\mu} (N_h - N_h^{eq}) - \frac{\Gamma_{N \rightarrow h}}{\mu} N_N = 0, \quad (25)$$

$$N_N' - \frac{\Gamma_{h \rightarrow N}}{\mu} (N_h - N_h^{eq}) + \frac{\Gamma_{N \rightarrow h}}{\mu} N_N = 0, \quad (26)$$

$$N_L' - \frac{\epsilon_h}{2} \frac{\Gamma_{h \rightarrow N}}{\mu} (N_h - N_h^{eq}) - \frac{\epsilon_N}{2} \frac{\Gamma_{N \rightarrow h}}{\mu} N_N = 0, \quad (27)$$

$$R' - \frac{\Gamma_{h \rightarrow f\bar{f}}}{\mu} \beta \varphi (N_h - N_h^{eq}) A - \frac{\Gamma_{h \rightarrow N}}{\mu} \beta \varphi (N_h - N_h^{eq}) A - \frac{\Gamma_{N \rightarrow h}}{\mu} \frac{g^{3/2}}{4\pi^3} \frac{m_N}{\dot{\phi}_0^{1/2}} N_N A = 0, \quad (28)$$

$$A' - A\mathcal{H} = 0, \quad (29)$$

where the prime superscript denotes derivative with respect to  $x$ , and  $\beta = g^{5/2} \dot{\phi}_0^{1/2} / (4\pi^3 \mu)$ .

The range of value used for the parameters during the numerical integration are as follows:  $3 \times 10^{-5} \text{ eV} < \tilde{m}_1 < 1 \text{ eV}$ ;  $10^9 \text{ GeV} < m_N < 10^{15} \text{ GeV}$ ;  $\mu > 10^{13} \text{ GeV}$ ;  $\dot{\phi}_0 < (10^{16} \text{ GeV})^2$ ; and  $g \sim 1$ . The upper limit to  $\tilde{m}_1$  comes from the sum of the three left-handed neutrino mass combining neutrino oscillation data [17] with constraints from the cosmic microwave background and large scale structure observations [16], under the assumption that  $\tilde{m}_1 \lesssim \sum m_\nu$  with a hierarchical left-handed neutrino spectrum. The lower limit has been arbitrarily set. The Yukawa coupling must be neither too small nor too large for the see-saw mechanism to be compelling. The upper bound of  $m_N$  is derived from Eq. (11) by setting  $(Y_\nu^\dagger Y_\nu)_{11} \sim 100$ . The inflation parameters  $\mu$  and  $\dot{\phi}_0$  are derived from observation [18]. In the preheating model we consider, the mass of the inflaton during preheating must be larger or equal to the mass of the inflaton during inflation. Hence we consider  $\mu > 10^{13} \text{ GeV}$ . We stress that the bounds of all the parameters are approximate and not very stringent.

We terminate preheating when  $\rho_R/\rho_\phi \geq 10$ , deeming this to be sufficient that the radiation energy dominates over the scalar energy density. An example of numerical results for a model is shown in Fig. 2.

Reheating is completed after twelve oscillations for this case, but only the first few oscillations are shown for the sake of clarity. It is clearly seen that  $h$  decays very quickly after it is created. Most of the decay happens in the  $h \rightarrow f\bar{f}$  channel. For efficient leptogenesis,  $\phi_c$  must be close to the origin in order for the  $h \rightarrow N$  process to start before the  $h$ s all decay away. Because of the quick decay of the  $h$ s, each oscillation can be treated as independent from each other. The decrease of the maximum point of  $\varphi$  indicates the energy of the  $\phi$  field going into thermalization of the universe.  $N_L$  is normalized with  $\epsilon(Y^\dagger Y)$  in the figure.

We have calculated the lepton number  $n_L/s$ , where  $s$  is the entropy. The lepton number gets translated into a baryon number  $n_B/s$  by sphaleron processes. Sphalerons transfer a lepton asymmetry to a baryon asymmetry by reactions conserving  $n_{B-L}$  but violating  $n_{B+L}$ . The relation between baryon number and lepton number is [19]

$$\frac{n_B}{s} = C \frac{n_{B-L}}{s} = \frac{C}{C-1} \frac{n_L}{s}; \quad C = \frac{24 + 4g_h}{66 + 13g_h}, \quad (30)$$

where  $g_h$  is the number of  $h$  generations. We consider a one-generation  $h$  model. The observation of baryon number comes from big bang nucleosynthesis (BBN) considerations [20], cosmic microwave background determinations [21], and large scale structure measurements [22] observations, with an assumed cosmological model. In a  $\Lambda$ CDM cosmology, these observations imply

$$\frac{n_B}{n_\gamma} = 6.5_{-0.2}^{+0.3} \times 10^{-10} \implies \frac{n_B}{s} \approx 9 \times 10^{-11} \quad (s = 7.04n_\gamma). \quad (31)$$

Figure 3 shows the region of  $m_N$  and  $\tilde{m}_1$  where  $n_B/s$  is higher than observation. We have used the maximally allowed CP parameter  $|\epsilon^{max}|$  in our Boltzmann equations, but  $|\epsilon|$  does not always retain the maximal value. The lower limit of  $\tilde{m}_1$  is due to the bound of  $3 \times 10^{-5} \text{ eV}$  we used in our calculations; if the bound is lowered, the contour simply continues downward. The slant shape on the left hand boarder of the shaded area is not a simple slope relation; this comes from the combined restriction of  $n_B/s \geq 9 \times 10^{-11}$  and the nonthermal condition of  $m_N > T$ . The preheat field parameters  $\mu$  and  $\dot{\phi}_0$  are not very sensitive in determining  $n_B/s$ .

The reheat temperature  $T_{RH}$  is greater than  $10^{10} \text{ GeV}$  in most of these regions. In supersymmetric models, this leads to overproduction of gravitinos which causes incompatibility with BBN observations [23]. Some models of supersymmetry have a larger mass to the gravitino [24], which can relax the constraint on  $T_{RH}$ . As we do not explicitly consider supersymmetry in our calculations above, our model agrees with all observations.

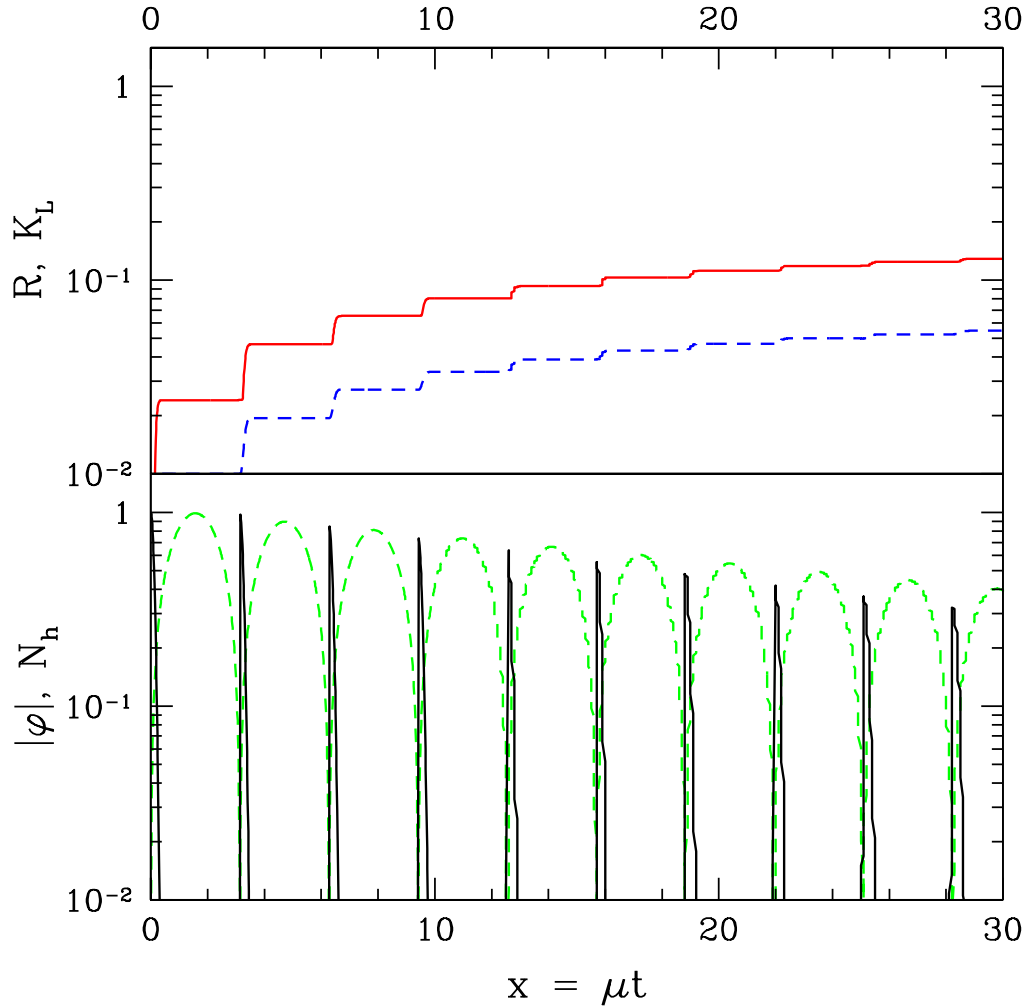


FIG. 2: Evolution of  $K_L \equiv N_L/(\epsilon(Y^\dagger Y)_{11})$  (solid line) and  $R$  (dashed line) in the upper box,  $|\phi|$  (dashed curve) and  $N_h$  (solid line) in the lower box, during preheating. Most of the lepton number is created during the early stage of oscillations. Note that  $h$  decays very quickly. Parameters used are  $\phi_0 = 10^{28} \text{ GeV}^2$ ,  $\mu = 10^{13} \text{ GeV}$ ,  $m_N = 10^{14} \text{ GeV}$ ,  $\tilde{m}_1 = 10^{-2} \text{ eV}$ , and  $g = 1$ .

#### IV. CONCLUSION

We have proposed a simple, economical model of nonthermal leptogenesis during instant preheating in the context of standard model and its extension to include Majorana partners. A hybrid inflation is employed, which allows us to evade the constraints on the properties of the inflaton potential from observations.

The assumed strong coupling between  $\phi$  and  $h$  ensures a quick thermalization. Preheating occurs within a few oscillations justifying the definition of “instant” in this preheating scenario. The dominant thermalization process is  $h \rightarrow f\bar{f}$ . In addition to this process, the see-saw mechanism provides another decay channel via the  $N$ - $h$  coupling. As  $m_h \propto \phi$ ,  $m_h$  can be larger or smaller than  $m_N$  throughout preheating. When  $m_h > m_N$ ,  $h \rightarrow N$  decay occurs, and the opposite happens when  $m_h < m_N$ . Both processes produce lepton number. The lepton asymmetry is produced under nonthermal conditions, since the  $T < m_N$ . The lepton number subsequently gets transformed to baryon number via sphaleron process. The effective light neutrino mass  $\tilde{m}_1$  has an upper bound of 1 eV. For a successful leptogenesis to happen in our model, we require  $m_N > 10^{11} \text{ GeV}$ . For most of the parameter space we find  $T_{RH} > 10^{10} \text{ GeV}$ , which may cause incompatibility with BBN observations in *some* SUSY models.

To summarize, if the electroweak Higgs is coupled to the inflaton, then one can expect instant preheating where the inflaton energy is extracted by resonant Higgs production as the inflaton passes through  $\phi = 0$ . As the inflaton grows during an oscillation, the effective mass of the Higgs may become large enough such that it can decay to the right-handed Majorana neutrino  $N$ , even if the mass of the  $N$  is as large as  $10^{11}$  to  $10^{16} \text{ GeV}$ . A lepton number may be produced in this phase. Later, when the value of the inflaton field decreases, the Higgs mass will fall below the  $N$

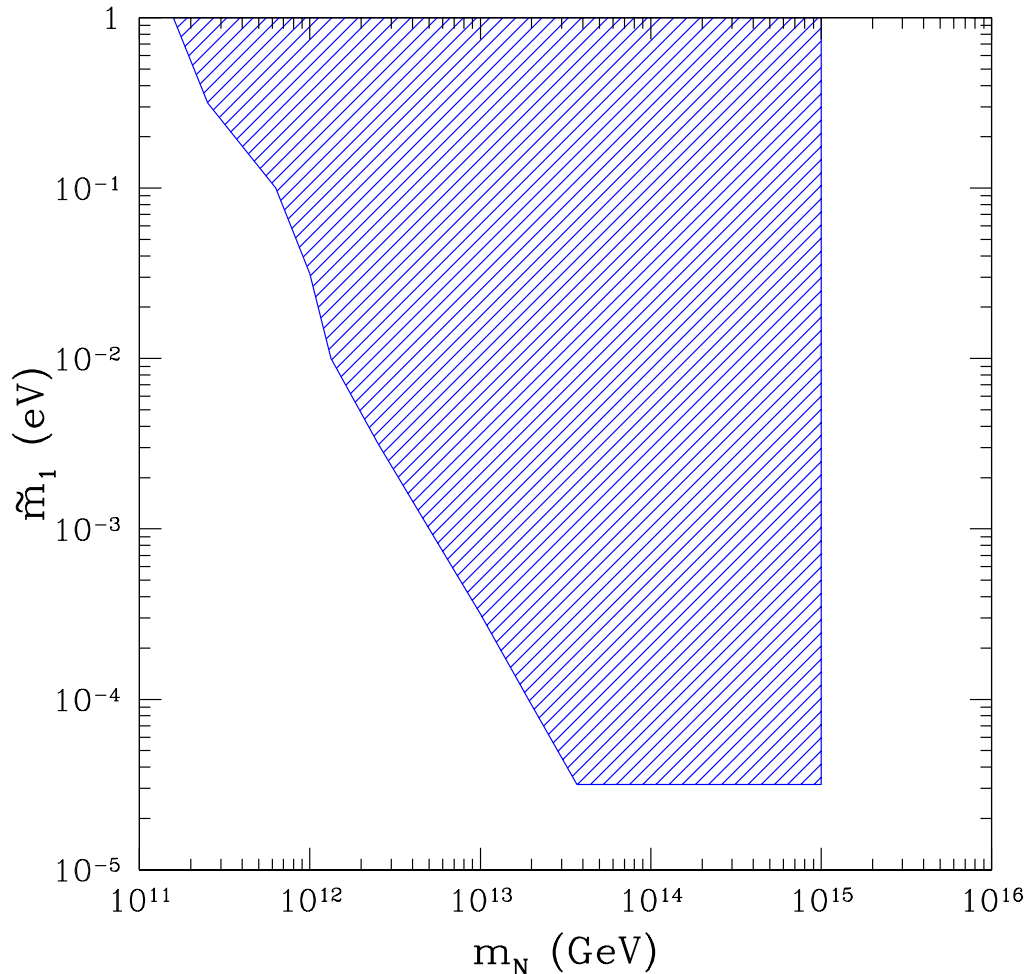


FIG. 3: Regions of  $m_N$  and  $\tilde{m}_1$  which satisfies the observed  $n_B/s$  value.

mass, and the  $N$  will decay to Higgs, also producing a lepton number.

The resulting lepton number only weakly depends on inflation parameters, is rather more sensitive to two neutrino mass parameters from the neutrino sector, and depends on the CP-phase in the heavy neutrino sector.

Our model is yet another scenario for leptogenesis, and illustrates the cosmological richness of the well motivated see-saw explanation for neutrino oscillations.

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